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Neutrino Mass Effects in a Minimally Extended Supersymmetric Standard Model

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Abstract

We consider an extension of the supersymmetric standard model which includes singlet Higgs superfield representations (in three generations) to generate neutrino masses via the see-saw mechanism. The resulting theory may then exhibit R-parity violation in the couplings of the singlets, inducing R-parity violating effective interactions among the standard model superfields, as well as inducing decay of the lightest neutralino, which otherwise would compose a stable LSP. We compute the rates for the resulting neutralino decays, depending on the particular superpotential couplings responsible for the violation of R-parity. We compare to astrophysical constraints on the decay of massive particles.

Despite its current experimental elusiveness, the supersymmetric standard model remains theoretically well motivated. Supersymmetry itself represents the unique possibility to combine internal and spacetime symmetries in quantum field theories, evading "no go" theorems by its incorporation of anticommutation relations in its defining algebraic structure. The gauging of supersymmetry inevitably results in a generally coordinate invariant theory, and Einsteinian gravity, and leads along the path of unification of gravity with the gauge interactions. Unification of the gauge interactions themselves appears to be facilitated by supersymmetry; extrapolation of the gauge coupling constants of the standard model gauge group factors according to the renormalization group equation with standard model matter content does not result in them coming together at a single point. On the other hand, when the superpartners of the standard model matter multiplets (including two Higgs doublets as required in the supersymmetric standard model) are included in the renormalization group running above the electroweak scale then the coupling constants for SU(3), SU(2), and U(1), cross at an energy scale of order 10^{16} GeV [1], as would be required for unification, and this unification scale is consistent, in these theories, with the observed stability of the proton. Finally, the inclusion of supersymmetry partners at about the electroweak scale is essential for the strongest phenomenological motivation for supersymmetry, which is to explain the stability of the electroweak scale under radiative corrections, and the maintenance of the hierarchy between the electroweak scale and the GUT or Planck scales.

It is well known, however, that the minimal supersymmetric standard model (MSSM) contains only a few of the possible gauge invariant couplings. With a minimal field content, the superpotential can be written as the sum of Yukawa terms (with generation indices suppressed)

$$F_Y = h_u H_1 Q u^c + h_d H_2 Q d^c + h_e H_2 L e^c \tag{1}$$

and to avoid an axion like state a mass term mixing the two Higgs doublets H_1 and H_2,

$$F_H = \mu H_1 H_2 \tag{2}$$

Q and L are weak doublets and u^c, d^c, and e^c are their corresponding right-handed counterparts. These are the only superpotential terms necessary to recover standard model fermion

masses and Higgs couplings. The MSSM possesses a Z_2 symmetry known as R-parity which can be represented by $R = (-1)^{F+L+3B}$ in terms of the particle's fermion, lepton, and baryon numbers.

One obvious extension of the MSSM consists of the inclusion of neutrino masses via a see-saw mechanism[2]. This is easily accomplished by the addition to the superpotential of

$$F_\nu = MNN + h_\nu H_1 LN \tag{3}$$

Neutrino masses of order ${h_\nu}^2 {v_1}^2/M$ will then be generated for the light (left-handed) neutrinos, where $v_1 = \langle H_1 \rangle$. It should be noted that the interactions induced by the superpotential F_ν do not violate R-parity as they they only violate lepton number in units of two. They do not, however, constitute the most general set of N-field interactions allowed by gauge invariance. As well as F_ν, one may also introduce the superpotential terms

$$F_N = \lambda H_1 H_2 N + kN^3 \tag{4}$$

The combination of the interactions in F_ν with *either* of the interactions in F_N will result in violation of R-parity.

The inclusion of neutrino mass by the see-saw mechanism has many other benefits in addition to the generation of neutrino masses which can in principle aid in the solution to the solar neutrino problem and/or atmospheric neutrino deficit and/or cosmological hot dark matter (though not all simultaneously without the inclusion of a fourth sterile neutrino). Right-handed neutrino decay has been utilized[3] to generate a lepton asymmetry which in conjunction with non-perturbative electroweak interactions becomes a baryon asymmetry. In refs. [4, 5], this mechanism was extended to supersymmetric models as well. Another possibility[6] for the generation of a baryon asymmetry made use of flat directions in the scalar potential as in the the Affleck-Dine mechanism[7]. In the latter, the superpotential $F = F_Y + F_N + F_\nu$ was required in order to induce a lepton number violating operator. For simplicity only one set (3 generations) of chiral superfields were added. Thus R-parity was explicitly violated. In that model, R-parity could have been preserved if the N-fields in F_N were distinct and have a different R-parity assignment than that of the N's in F_ν.

There are numerous other ways in which one can imagine extending the MSSM. In what is often called the minimal-nonminimal superymmetric standard model (MNMSSM) a single additional gauge singlet chiral superfield, N is added[8]. This extension is realized by simply adding to the superpotential the contribution from F_N (4). The primary motivation for the inclusion of the Higgs singlet is the possibility that it offers for the dynamical generation of the Higgs mixing mass μ. If the N field is a field which acquires a vev determined by mass parameters of the order of the electroweak scale, then with a NH_1H_2 coupling of standard strength (say comparable to a gauge coupling) Higgs mixing of the requisite magnitude is induced. On the other hand, if the mass parameters in the N sector are much larger, say of an intermediate scale, or perhaps of the GUT scale, as might naturally be expected to be in see-saw models, then if the N has a nonzero vev one would naturally expect it to also be of this scale. In such a case one still might imagine inducing a weak scale mixing between the Higgs doublets, at the price of fine tuning the NH_1H_2 coupling to be small to give the hierarchical ratio between the electroweak scale and the N mass scale. Though this small (O(M_W) mixing mass is technically feasible its smallness is part and parcel with the hierarchy problem. The cubic term is required in order to avoid an N-axion like field, in the absence of an explicit μ superpotential term mixing the two Higgs supermultiplets. In a detailed exaimation of this model[9], it was found that many of the standard Higgs mass relations are altered. If the MSSM Higgs mass relations are found to be experimentally not viable, this model becomes the simplest alternative.

From another point of view, the MNMSSM is of interest as it can easily produce a relatively light dark matter candidate [10]. In the MSSM, steadily improving accelerator limits, are pushing up the mass of the lightest supersymmetric particle (LSP), which due to the unbroken R-parity in the MSSM is stable. In the minimal model the LSP is generally expected to be a linear combination of the four neutral $R = -1$ fermions[11], the two gauginos, $\tilde{B}$ and $\tilde{W}$, and the Higgsinos $\tilde{H}_1$ and $\tilde{H}_2$. With regards to a dark matter candidate, the best choice in the MSSM appears to be the bino whose mass is typically between 40 GeV and $\sim$ 300 GeV for cosmologically interesting parameters[12]. In the non-minimal model it is quite feasible[10, 13] to have a a light LSP (10 - 50 GeV), which is a state which has a strong admixture the fernionic component of N. Though, cosmologically, a very massive

LSP is just as good as a light one (light still referring to O(GeV)), from the point of view of experimental detection, the lighter one is better[14].

In this letter we derive the consequences of the R-parity violation of the full superpotential. R-parity violation in the quark sector is usually avoided in order to insure a relatively stable proton. In the Higgs-lepton sector, there are many constraints on R-parity violation as well. In the case we consider here, R-parity is violated only in the heavy N-field sector. Nevertheless, this R-parity violation shows up in the low energy sector, most notably in the destabilization of the LSP. We derive constraints on the neutrino mass parameters as a consequence of the constraints on late-decaying LSP's.

As well as the destabilization of the LSP to which we will turn below, there are other possible low-energy signatures of R-parity violation in the high energy N-field sector. If supersymmetry were exact, then even the combined presence of the F_ν and F_N superpotential terms would not induce (super)renormalizable lepton number violating superpotential terms involving only the light superfields of the theory, due to the nonrenormalization theorems for the superpotential. After supersymmetry breaking the nonrenormalization theorems no longer hold exactly, and lepton number (and hence R-parity)violating effective interactions will be induced in an amount governed by the scale of supersymmetry breaking. This will result in low energy R-parity violating interactions involving standard model superfields of the form of both induced effective superpotential terms such as

$$F_{RX} = m_X H_1 L + \lambda_X L L e^c \tag{5}$$

as well as soft supersymmetry breaking lepton number violating terms. By appropriate change of basis we may diagonalize the Higgs-lepton mass mixing and parametrize our lepton number violating effects by λ_X. These terms are induced from one loop diagrams in amounts

$$\lambda_X \sim \frac{m_\delta^2}{M_N^2} \qquad or \qquad \lambda_X \sim \frac{\mu m_\delta}{M_N^2} \tag{6}$$

where m_δ is the scale of supersymmetry breaking. Lepton number violating renormalizable interactions of this type are constrained by laboratory limits on lepton flavour violation, and neutrinoless double beta decay [15]. As we have analyzed previously, even stronger

limits are imposed on interactions of this type by the persistence of a baryon asymmetry in the early universe, assuming that it is not produced at or after the electroweak phase transition [16, 17]. The danger here is that the lepton number violation implied by the new interaction could attain thermal equilibrium at the same time as baryon and lepton number violating (but B-L conserving) nonperturbative electroweak interaction effects to simultaneously equilibrate both the baryon and lepton number of the universe to zero. If these limits pertain, they would imply that $\lambda_X < 7 \times 10^{-7}$[16]. These limits may be evaded, and indeed a baryon asymmetry may be generated from a lepton asymmetry, provided one of the generations of lepton flavours has its lepton number violating interaction in equilibrium, while another does not [18].

As we have mentioned above, the combination of the NH_1L superpotential term with either the NH_1H_2 or NNN superpotential interactions breaks R-parity and hence will destabilize the lightest neutralino mass eigenstate. The nature of the resulting decay will depend on which of these latter terms is responsible. Let us begin our discussion with consideration of decays induced by the NH_1H_2 term. There will be tree-level two body decays to lepton-Higgs final states induced by the diagrams shown in figure 1(a), and 1(b). We see that for decay from an $\tilde{H}_2$ component of a neutralino, the vev is the large H_1 vev, favoring that amplitude over the amplitude for the decay from diagram 1(a) with decay from the $\tilde{H}_1$ component of the neutralino, by a factor of $\tan\beta = v_1/v_2$, the ratio of the vevs. In addition there will be favoured (by $\tan\beta$) decay amplitudes for the decays from the $\tilde{H}_2$ and $\tilde{H}_1$ components of the neutralino coming from figure 1(b). To get the approximate contribution to the decay of the amplitude of figure 1(a), we note that the insertion of the Higgs vev induces a mass mixing, of the neutralino with the N field, of magnitude m_{NHH}/M, where m_{NHH} is the mixing mass $m_{NHH} = \lambda v \sin\beta$, where $v^2 = v_1^2 + v_2^2$ and M is the Majorana mass term for the N-field. The N-field component of the resulting mass eigenstate then induces a decay to an H_1L final state (suppressing lepton generation indices) via the $h_\nu NH_1L$ superpotential coupling. Similarly, in diagram 1(b), insertion of the Higgs vev induces a mass mixing of the outgoing lepton with the N-field, of magnitude m_ν/M, where m_ν is the Dirac neutrino mass $m_\nu = h_\nu v \sin\beta$, and M is the Majorana mass term for the N-field. These mixings then appear in decay amplitudes induced by the coupling at the other vertex, into a two body

final state, with decay width (ignoring mixing factors) of the form:

$$\Gamma_o \simeq \frac{m_{\tilde{\chi}_o}}{16\pi}(1 - \frac{m_o^2}{m_{\tilde{\chi}_o}^2}) \tag{7}$$

where m_o is the mass of the final state Higgs scalar. If (in the absence of mixing with the N-field) we would write the LSP as an admixture

$$\tilde{\chi}_o = \alpha\tilde{B}^o + \beta\tilde{W}^o + \gamma\tilde{H}_1^o + \delta\tilde{H}_2^o \tag{8}$$

Then the decay of the LSP via its H_2 component will then occur at a rate

$$\Gamma \simeq \delta^2 \frac{4\lambda^2 h_\nu^2 v^2 \sin^2\beta}{M_N^2}\Gamma_o \tag{9}$$

while there would be a contribution to the neutralino decay width from decay of its H_1 component with a contribution

$$\Gamma \simeq \gamma^2 \frac{\lambda^2 h_\nu^2 v^2 \sin^2\beta}{M_N^2}\Gamma_o \tag{10}$$

There will also be decays into neutrino-gamma modes induced at one loop, as shown in figure 2. They will give a decay rate with the same mixing factors as the tree level modes, multiplied by loop induced dipole decay width.

Now to produce a two body decay to a neutrino plus physical (on mass shell) photon, the only part of the electromagnetic vertex which contributes is the induced transition dipole piece, which we may parametrize as [19]

$$M_\mu = -i\bar{u}(p_f)\frac{i\sigma_{\mu\nu}q^\nu}{(m_{\tilde{\chi}_o} + m_\nu)}(F_2^V + F_2^A\gamma_5)u(p_i) \tag{11}$$

which results in a dipole decay rate (dropping the neutrino mass)

$$\Gamma_D = \frac{m_{\tilde{\chi}_o}}{8\pi}\left[[F_2^V]^2 + [F_2^A]^2\right] \tag{12}$$

From inspection of the diagram we find that the mixing factors associated with the N mass and the Higgs vevs must combine with the kinematics of the dipole decay to give a net decay width of order

$$\Gamma \simeq \frac{m_{\tilde{\chi}_o}}{8\pi}\frac{v^4 \sin^4\beta h_\nu^2 \lambda^2 e^2}{M_N^2 m^2} \tag{13}$$

where m is a mass of the order the electroweak scale. Note that in order to induce a dipole matrix element we have to have broken supersymmetry. This implicitly appears in our estimate in that lines in the loop, which is dominated by momenta of order the electroweak scale, are split in mass by supersymmetry breaking of order that scale, giving a result whose magnitude we may read off an individual diagram as above. We also note that there is no diagram involving the N-field in a loop in such a way as to induce a dipole with less suppression by the N-field mass, as such diagrams involve photon emission from external lines, and the Ward-Takahashi identities of electromagnetic gauge invariance ensure that such terms do not contribute to the induced dipole [19].

Similarly, decays of the LSP may be induced by the NNN superpotential term, as represented by the diagrams of figure 3. We note that figure 3(a) is an induced D-term and contains a loop which is also a D-term. This ensures a non-zero decay rate for the neutralino even when supersymmetry is unbroken, unlike the case for F-terms. Because D-terms do not obey non-renormalization theorems, they can be radiatively induced even in the limit of unbroken supersymmetry; hence in general they appear without suppression factors associated with the scale of supersymmetry breaking.We also note that the induced D-term in figure 3(a) (and its associated component diagrams) is a dimension six term [20]. The component diagrams relevant to neutralino decay are shown in figure 3(b) to 3(e). The processes of figure 3 dominate over decays induced by tree-level diagrams for large M_N, as the latter are suppressed by eight powers of M_N in rate, whereas the loop induced decays are only suppressed by four powers. Computing the diagrams of figures 3(b) and 3(c) one finds that they result in a decay rate that is approximately

$$\Gamma \sim \gamma^2 \frac{k^2 h_\nu^6}{16\pi(2\pi)^8} \frac{\mu^2 v_1^2 m_{\tilde{\chi}_o}}{M_N^4} \tag{14}$$

whereas the final two diagrams of figure 3 result in a decay rate for the LSP that is approximately

$$\Gamma \sim \gamma^2 \frac{k^2 h_\nu^6}{16\pi(2\pi)^8} \frac{v_1^2 m_{\tilde{\chi}_o}^3}{M_N^4} \tag{15}$$

We expect that the Higgsino mass should be at least of the order of the doublet mixing term, and in certain circumstances the doublet mixing term might be substantially smaller;

so below we will use the latter of these rate estimates for numerical estimates.

Almost without exception, the LSP decays we are considering are effectively entropy producing decays, ie. they will produce high energy photons. Photon producing decays are known to be highly constrained from both astrophyical and cosmological observations (see eg. ref. [21] for a recent compilation of such limits). These limits generally place constraints in the density-lifetime plane of the decaying particle. We will assume that the LSP, in the absense of its decay, is the dominant form of dark matter and therefore assume that its cosmological density is such that $\Omega_\chi \approx 1$, where $\Omega = \rho/\rho_c$ is the cosmological density parameter. At this density, one finds that the LSP lifetime is constrained so that either $\tau_\chi \lesssim 10^4$s to avoid affecting the light element abundances produced during big bang nucleosynthesis, or the LSP must be effectively stable with a lifetime $\tau_\chi \gtrsim 10^{24}$s. Astophysical limits on other R- parity violaing interactions were considered in [22].

The decay rates in Eqs. (9,10,15) are clearly dependent on a number of model parameters. In order to get a feeling for the limits imposed by the cosmological constraints we make a few more assumptions. We assume that the LSP is primarily a gaugino (a bino) with mass $m_\chi \approx 150$ GeV. For $|\mu| \sim 1 - 10$ TeV, $\gamma \sim 2 \times 10^{-3} - 2 \times 10^{-2}$ and $\delta \sim 4 \times 10^{-3} - 4 \times 10^{-2}$ and for large $\tan\beta$, $\sin\beta \approx 1$. We can then write (for the decays based on the $H_1 H_2 N$ superpotential term)

$$\tau_\chi \simeq 3 \times 10^{-6} h_\nu^{-2} \lambda^{-2} (4\delta^2 + \gamma^2)^{-1} \left(\frac{M_N}{10^{12}\mathrm{GeV}}\right)^2 \left(\frac{150\mathrm{GeV}}{m_\chi}\right) \mathrm{s} \tag{16}$$

Taking central values for γ and δ, and $h_\nu \sim \lambda \sim h$, we have

$$\tau_\chi \simeq 7 \times 10^{-3} h^{-4} \left(\frac{M_N}{10^{12}\mathrm{GeV}}\right)^2 \left(\frac{150\mathrm{GeV}}{m_\chi}\right) \mathrm{s} \tag{17}$$

The constraints on τ_χ are therefore constraints on M_N,

$$M_N \lesssim 10^{15} h^2 \ \mathrm{GeV} \tag{18}$$

or

$$M_N \gtrsim 10^{25} h^2 \ \mathrm{GeV} \tag{19}$$

The latter limit, of course only makes sense for $h \ll 1$ and in this case the LSP is effectively stable as its lifetime is much greater than the age of the Universe.

For LSP decay induced by the kN^3 superpotential term, from the decay width estimate given above we deduce an LSP lifetime of order

$$\tau_\chi \simeq 4 \times 10^{20} {h_\nu}^{-6} k^{-2} \gamma^{-2} \left(\frac{M_N}{10^{12}\mathrm{GeV}}\right)^4 \left(\frac{150\mathrm{GeV}}{m_\chi}\right)^3 \mathrm{s} \tag{20}$$

which translates into the limits

$$M_N \lesssim 5 \times 10^6 {h_\nu}^{3/2} k^{1/2} \ \mathrm{GeV} \tag{21}$$

or

$$M_N \gtrsim 5 \times 10^{11} {h_\nu}^{3/2} k^{1/2} \ \mathrm{GeV} \tag{22}$$

These limits show therefore that even if R-parity violation is inserted in the singlet sector, destabilization of the LSP can indeed occur and R-parity violation of this type is strongly constrained. It is especially interesting that cosmological arguments provide such strong constraints, probing possible see-saw sources of R-parity violation to far higher mass scales than could be directly accessed by laboratory experiment; this provides yet another example of the power of cosmological considerations to provide us with new information about the fundamental interactions of nature.

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Figure Captions

Figure 1: Neutralino decay diagrams induced by an NH_1H_2 superpotential term.

Figure 2: Radiative neutralino decay diagrams induced by an NH_1H_2 superpotential term.

Figure 3: Neutralino decay diagrams induced by an NNN superpotential term. Figure 3(a) is the superfield diagram whose dominant associated component field diagrams include those shown in Figures 3(b) through 3(e).